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Effects of Magnetic and Collisional Viscosity on the E X B Gradient Drift Instability

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We investigate the role of magnetic viscosity (i.e., finite Larmor radius effects) and collisional viscosity on the E x B gradient drift instability in both the collisional and inertial regimes. We derive the equations describing the time evolution of small perturbations to an equilibrium two dimensional (x,y) magnetized plasma, where the equilibrium electron density is an arbitrary function of y. The equations are then solved iteratively in time to obtain the growth rate of the fastest growing eigenmode for a given Fourier wavenumber k in the x direction. We compare our results for appropriate profiles with two asymptotic results: the long wavelength limit, valid for $kL << 1$, and the short wavelength or local limit, valid for $kL >> 1$, where L is the gradient scale length of the equilibrium profile. It is found that although the long and short wavelength limits do indeed provide accurate growth rates for $kL << 1$ and $kL >> 1$, respectively, neither provides reliable growth rates for $kL \sim 1$. In comparing our results to the long wavelength limit, we conclude that in

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addition to affecting γ in an absolute sense, increasing L to a finite value decreases k_{max}, the value of k for which $\gamma(k)$ maximizes, and also affects the sharpness of the peak of $\gamma(k)$ about k_{max}. Both of these phenomena have large effects on freezing models.

CONTENTS

I.	INTRODUCTION	1
II.	FUNDAMENTAL EQUATIONS	2
III.	STABILITY ANALYSIS	9
IV.	ASYMPTOTIC RESULTS: SHORT- AND LONG-WAVELENGTH LIMITS	11
	A. SHORT-WAVELENGTH LIMIT	11
	B. LONG-WAVELENGTH LIMIT	15
V.	NUMERICAL RESULTS	16
	A. LONG-WAVELENGTH COMPARISONS	16
	B. SHORT-WAVELENGTH COMPARISONS	18
	C. RELATIVE EFFECTIVENESS OF η_1 AND η_3	19
VI.	CONCLUSIONS	19
	ACKNOWLEDGMENT	20
	REFERENCES	21
	DISTRIBUTION LIST	33

EFFECTS OF MAGNETIC AND COLLISIONAL VISCOSITY ON THE $\underline{E}$ X $\underline{B}$ GRADIENT DRIFT INSTABILITY

I. INTRODUCTION

Much has been written in the recent past regarding the 'freezing' phenomenon exhibited by ionospheric ion clouds. Recently the debate has centered upon whether the explanation for the phenomenon is a 2- or 3-dimensional effect. We shall not enter that debate here, but rather investigate one of the more promising 2-D effects, proposed by Sperling and Glassman (1985). They propose that collisional and magnetic viscosities, η_1 and η_3, respectively, [Braginskii, 1965] suppress the growth of the gradient drift instability at short wavelengths, and that this effect can explain the 'freezing' phenomenon in both barium ion clouds and nuclear plumes. The nonlocal stability analysis based on the full stress tensor presented in Sperling and Glassman (1985) considered only the asymptotic case of $kL \ll 1$ where k is a Fourier wavenumber and L is the gradient scale length. The results presented in their paper were quite intriguing and showed that η_1 and η_3 can exhibit a strong stabilizing effect on the gradient drift instability.

However, one must seriously question the applicability of the above work to barium clouds or nuclear striations since the results are valid only in the long wavelength limit which corresponds to $L \to 0$, and the stress tensor itself is valid only for $L \gg \rho_i$ where ρ_i is the ion gyroradius. What are the instability characteristics for <u>finite</u> L, and in particular for $kL \sim O(1)$ as we expect it to be for some cases of interest. This poses a difficult numerical problem, and has not henceforth yielded to solution. However, we have recently developed numerical techniques which allow us to solve the linearized equations for the full stress tensor, and to investigate the applicability and validity of the long wavelength asymptotic results. We find that the gradient scale length L does have a strong effect on the growth rate, and thus that long

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wavelength analysis is often inadequate for realistic parameters.

II. FUNDAMENTAL EQUATIONS

The equations describing a two dimensional plasma in a plane perpendicular to an ambient magnetic field $\underset{\sim}{B}$ have been given in many places [e.g., Mitchell et al., 1985]. What is new here is the presence of terms which account for magnetic and collisional viscosity. As with Sperling and Glassman (1985), we shall ignore all electron collisions. We shall do this because we do not wish to address the effects of these collisions (primarily diffusion) here, and because the diffusion induced by these collisions precludes the attainment of a true equilibrium upon which to perform a stability analysis.

Consider a two dimensional plasma consisting of ions and electrons, embedded in a neutral gas and in a constant ambient magnetic field $\underset{\sim}{B}$ perpendicular to the plasma plane. If we neglect ion-electron collisions, the momentum equation describing species α is

$$\left(\frac{\partial}{\partial t} + \underset{\sim}{v}_\alpha \cdot \nabla\right)\underset{\sim}{v}_\alpha = \frac{q_\alpha}{m_\alpha}\left(\underset{\sim}{E} + \frac{\underset{\sim}{v}_\alpha \times \underset{\sim}{B}}{c}\right) - \nu_{\alpha n}\left(\underset{\sim}{v}_\alpha - \underset{\sim}{U}_n\right) - \frac{\nabla P_\alpha}{n_\alpha m_\alpha} + \underset{\sim}{g} + \frac{\underset{\sim}{L}_\alpha}{m_\alpha} \qquad (1)$$

where the subscript α denotes the plasma species (i for ions, e for electrons, for example), n is the species number density, $\underset{\sim}{v}$ is the species fluid velocity, $P_\alpha = n_\alpha k_B T_\alpha$ is pressure, $\underset{\sim}{E}$ is the electric field, $\underset{\sim}{g}$ is the gravitational acceleration, q is the species charge, $\nu_{\alpha n}$ is the species collision frequency with the neutral gas, $\underset{\sim}{U}_n$ is the neutral wind velocity, c is the speed of light, k_B is Boltzmann's constant, m is the species particle mass, and $\underset{\sim}{L}$ is the "force" per particle due to the η_1 and η_3 terms in the Braginskii stress tensor. We can rewrite this equation as

$$\underset{\sim}{F}_\alpha/m_\alpha + \frac{q_\alpha}{m_\alpha c}\left(\underset{\sim}{v}_\alpha \times \underset{\sim}{B}\right) = 0 \tag{2}$$

where

$$\underset{\sim}{F}_\alpha = \underset{\sim}{F}_{\alpha 1} + \underset{\sim}{F}_{\alpha 2}(\underset{\sim}{v}_\alpha) \tag{3}$$

$$\underset{\sim}{F}_{\alpha 1} \equiv q_\alpha \underset{\sim}{E} + m_\alpha \underset{\sim}{g} + \nu_{\alpha n} m_\alpha \underset{\sim}{U}_n - \nabla P_\alpha/n_\alpha \tag{4}$$

$$\underset{\sim}{F}_{\alpha 2}(\underset{\sim}{v}_\alpha) \equiv -\left(\frac{\partial}{\partial t} + \underset{\sim}{v}_\alpha \cdot \nabla\right)\underset{\sim}{v}_\alpha m_\alpha + \underset{\sim}{L}_\alpha(\underset{\sim}{v}_\alpha) - \nu_{\alpha n} m_\alpha \underset{\sim}{v}_\alpha \tag{5}$$

If we place ourselves in a Cartesian coordinate system in which $\underset{\sim}{B}$ is aligned along the z axis then (2) yields

$$\underset{\sim}{v}_\alpha = \frac{c}{q_\alpha B}\,\underset{\sim}{F}_\alpha \times \hat{e}_z. \tag{6}$$

Strictly speaking, this equation only applies to the perpendicular component of $\underset{\sim}{v}_a$. We assume $F_{\alpha||} = 0$ so that $v_{\alpha||} = 0$. Note that (6) above is actually an implicit expression for $\underset{\sim}{v}_\alpha$, since $\underset{\sim}{F}_{\alpha 2}$ is a function of $\underset{\sim}{v}_\alpha$. Hence $\underset{\sim}{v}_\alpha$ is solved iteratively. We define

$$\underset{\sim}{v}_{\alpha 1} = \frac{c}{q_\alpha B}\,\underset{\sim}{F}_{\alpha 1} \times \hat{e}_z \tag{7}$$

and

$$\underset{\sim}{v}_{\alpha 2} = \frac{c}{q_\alpha B}\,\underset{\sim}{F}_{\alpha 2}(\underset{\sim}{v}_{\alpha 1}) \times \hat{e}_z\ . \tag{8}$$

Then $\underset{\sim}{v}_\alpha$ can be approximated as

$$\underset{\sim}{v}_\alpha = \underset{\sim}{v}_{\alpha 1} + \underset{\sim}{v}_{\alpha 2} \tag{9}$$

The success of this procedure obviously depends on $\underset{\sim}{F}_{\alpha 2}$ being small with respect to $\underset{\sim}{F}_{\alpha 1}$, i.e., $\underset{\sim}{v}_{\alpha 2} \ll \underset{\sim}{v}_{\alpha 1}$.

The reader will perhaps be disturbed by the absence of explicit collision-frequency-dependent Pedersen and Hall mobilities. Indeed, the above procedure is not the one usually followed, wherein the term $\nu_{\alpha n} m_\alpha \underset{\sim}{v}_\alpha$ is included in (2) rather than (5) (e.g., Zalesak et al., 1985). The procedure used here is less accurate than this usual approach (e.g., Hall currents due to electric fields are totally absent here). However, we use this simplified set here for three reasons. First, it is accurate as long as $\nu_{in}/\Omega_i \ll 1$ and $|F_{\alpha 2}| \ll |F_{\alpha 1}|$; second, it is simple; and third, it is the approach used by Sperling and Glassman (1985), with whom we wish to compare our results.

We now list the additional assumptions we need to make in order to recover the linearized equations of Sperling and Glassman (1985). We assume that both the ion temperature T_i and the electron temperature T_e are constants in space and time, neglect all electron collisions, neglect gravity $\underset{\sim}{g}$, assume quasi-neutrality ($n_i = n_e = n$), assume singly changed ions ($q_i = -q_e = e$), and assume that the electric field is electrostatic (i.e., $\underset{\sim}{E} = -\nabla\phi$). The smallness of m_e, together with our neglect of electron collisions, allows us to neglect all but the first and third terms on the right-hand side of (1) for electrons. Thus

$$\underset{\sim}{v}_e = \underset{\sim}{v}_{e1} = \frac{c}{B}\left[\underset{\sim}{E} + (k_B T_e/e)(\nabla n/n)\right] \times \hat{e}_z \qquad (10)$$

since $\underset{\sim}{F}_{e2} = 0$.

The ion velocities are given by

$$\underset{\sim}{v}_{i1} = \frac{c}{eB}\left[e\underset{\sim}{E} + \nu_{in} m_i \underset{\sim}{U}_n - k_B T_i \nabla n/n\right] \times \hat{e}_z \qquad (11)$$

$$\underset{\sim}{v}_{i2} = \frac{c}{eB}\left[-\left(\frac{\partial}{\partial t} + \underset{\sim}{v}_{i1} \cdot \nabla\right)\underset{\sim}{v}_{i1} m_i + L_i(\underset{\sim}{v}_{i1}) - \nu_{in} m_i \underset{\sim}{v}_{i1}\right] \times \hat{e}_z \qquad (12)$$

$$\underset{\sim}{v}_i = \underset{\sim}{v}_{i1} + \underset{\sim}{v}_{i2} \tag{13}$$

Letting

$$\Phi = \phi - (k_B T_e/e) \ln n \tag{14}$$

$$\psi = \phi + (k_B T_i/e) \ln n \tag{15}$$

we get

$$\underset{\sim}{v}_e = - \frac{c}{B} \left(\nabla \Phi \times \hat{e}_z \right) \tag{16}$$

$$\underset{\sim}{v}_{i1} = - \frac{c}{B} \left(\nabla \psi - \frac{\nu_{in} m_i}{e} \underset{\sim}{U}_n \right) \times \hat{e}_z \tag{17}$$

To solve for the evolution of the plasma, we shall need two equations: (1) a continuity equation for electrons or ions, and (2) an equation for current continuity and quasi-neutrality, i.e., $\nabla \cdot \underset{\sim}{J} = 0$ where $\underset{\sim}{J}$ is the electric current $en(\underset{\sim}{v}_i - \underset{\sim}{v}_e)$. We choose the electron continuity equation for simplicity:

$$\frac{\partial n}{\partial t} = - \nabla \cdot (n \underset{\sim}{v}_e) = \nabla \cdot \left(\frac{c}{B} n \nabla \Phi \times \hat{e}_z \right) \tag{18}$$

The current density is given by

$$\underset{\sim}{J} = ne \left(\underset{\sim}{v}_i - \underset{\sim}{v}_e \right) = ne \left(\underset{\sim}{v}_{i2} + \frac{\nu_{in}}{\Omega_i} \underset{\sim}{U}_n \times \hat{e}_z - \frac{c}{B} k_B \left(T_i + T_e \right) \frac{\nabla n}{n} \times \hat{e}_z \right) \tag{19}$$

thus

$$\nabla \cdot J = \nabla \cdot \left[ne\left(\underset{\sim}{v}_{i2} + \frac{\nu_{in}}{\Omega_i} \underset{\sim}{U}_n \times \hat{e}_z\right)\right] \tag{20}$$

If we denote the x and y components of $\underset{\sim}{v}_{i1}$ as u and v, respectively, and denote partial differentiation by subscripts, then we can write

$$\underset{\sim}{L}(v_{i1}) = \frac{1}{n}\left[\left(\frac{\eta_0}{3} D\right)_x + \left(\eta_1 M\right)_x + \left(\eta_3 N\right)_x + \left(\eta_1 N\right)_y - \left(\eta_3 M\right)_y\right]\hat{e}_x$$

$$+ \frac{1}{n}\left[\left(\frac{\eta_0}{3} D\right)_y - \left(\eta_1 M\right)_y - \left(\eta_3 N\right)_y + \left(\eta_1 N\right)_x - \left(\eta_3 M\right)_x\right]\hat{e}_y \tag{21}$$

where

$$D = u_x + v_y = 0 \tag{22}$$

$$M = u_x - v_y = -2\frac{c}{B}\left(\psi_{xy}\right) \tag{23}$$

$$N = u_y + v_x = \frac{c}{B}\left(\psi_{xx} - \psi_{yy}\right) \tag{24}$$

$$\eta_0 = 0.96\ \bar{\eta} \tag{25}$$

$$\eta_1 = \bar{\eta}\ \frac{1.2\lambda^2 + 2.23}{\lambda^4 + 4.03\lambda^2 + 2.33} \tag{26}$$

$$\eta_3 = \bar{\eta}\ \frac{\lambda^3 + 2.38\lambda}{\lambda^4 + 4.03\lambda^2 + 2.33} \tag{27}$$

$$\bar{\eta} = \frac{n_i K_B T_i}{\nu_{ii}} \tag{28}$$

$$\lambda = \frac{2\Omega_i}{\nu_{ii}} \tag{29}$$

$$\Omega_i = \frac{eB}{m_i c} \tag{30}$$

$$\nu_{ii} = \frac{23.4 - 1.15 \log_{10}(n_i) + 3.45 \log_{10}(T_i)}{3 \times 10^7} \left(\frac{2}{A}\right)^{1/2} \frac{n_i}{T_i^{3/2}} \tag{31}$$

Here ν_{ii} is the ion-ion collision frequency, A is the ion atomic mass, and Ω_i is the ion cyclotron frequency [Braginskii, 1965, Sperling and Glassman, 1985]. T_i in (31) is expressed in eV.

Noting that

$$\underset{\sim}{v}_{i1} \times \hat{e}_z = \frac{c}{B} \left(\nabla\psi - \frac{\nu_{in} m_i}{e} \underset{\sim}{U}_n\right) \tag{32}$$

we obtain equations

$$\underset{\sim}{v}_{i2} = \frac{c}{eB} \left[- m_i \left(\frac{\partial}{\partial t} + \underset{\sim}{v}_{i1} \cdot \nabla + \nu_{in} \right) \frac{c}{B} \left(\nabla\psi - \frac{\nu_{in} m_i}{e} \tilde{U}_n\right)\right]$$

$$+ \frac{c}{eB} \underset{\sim}{L}_i(\underset{\sim}{v}_{i1}) \times \hat{e}_z \tag{33}$$

and

$$\nabla \cdot \underset{\sim}{J} = \nabla \cdot \left[ne \frac{\nu_{in}}{\Omega_i} \underset{\sim}{U}_n \times \hat{e}_z\right] + \nabla \cdot \underset{\sim}{J}_v$$

$$- \nabla \cdot \left[\frac{c^2}{B^2} n m_i \left(\frac{\partial}{\partial t} + \underset{\sim}{v}_{i1} \cdot \nabla + \nu_{in}\right)\left(\nabla\psi - \frac{\nu_{in} m_i}{e} \underset{\sim}{U}_n\right)\right] \tag{34}$$

where

$$\nabla \cdot \underset{\sim}{J}_v = \nabla \cdot \left(n \frac{c}{B} \underset{\sim}{L}_i \left(\underset{\sim}{v}_{i1}\right) \times \hat{e}_z\right) = \nabla \cdot \underset{\sim}{J}_Q + \nabla \cdot \underset{\sim}{J}_R \tag{35}$$

$$\nabla \cdot \underset{\sim}{J}_Q = \frac{c}{B} \left\{ \frac{\partial}{\partial x} \left[- \left(\eta_1 M\right)_y + \left(\eta_1 N\right)_x\right] - \frac{\partial}{\partial y} \left[\left(\eta_1 M\right)_x + \left(\eta_1 N\right)_y\right]\right\} \tag{36}$$

$$= \frac{c^2}{B^2}\left\{\frac{\partial}{\partial x}\left[\eta_1\left(\psi_{xxx}+\psi_{xyy}\right)+2\eta_{1y}\psi_{xy}+\eta_{1x}\left(\psi_{xx}-\psi_{yy}\right)\right]\right.$$

$$\left.-\frac{\partial}{\partial y}\left[-\eta_1\left(\psi_{yyy}+\psi_{xxy}\right)-2\eta_{1x}\psi_{xy}+\eta_{1y}\left(\psi_{xx}-\psi_{yy}\right)\right]\right\} \qquad (37)$$

$$\nabla\cdot\underset{\sim}{J}_R = \frac{c}{B}\left\{\frac{\partial}{\partial x}\left[-\left(\eta_3 N\right)_y-\left(\eta_3 M\right)_x\right]-\frac{\partial}{\partial y}\left[\left(\eta_3 N\right)_x-\left(\eta_3 M\right)_y\right]\right\} \qquad (38)$$

$$= \frac{c^2}{B^2}\left\{\frac{\partial}{\partial x}\left[\eta_3\left(\psi_{yyy}+\psi_{xxy}\right)+2\eta_{3x}\psi_{xy}-\eta_{3y}\left(\psi_{xx}-\psi_{yy}\right)\right]\right.$$

$$\left.-\frac{\partial}{\partial y}\left[\eta_3\left(\psi_{xxx}+\psi_{xyy}\right)+2\eta_{3y}\psi_{xy}+\eta_{3x}\left(\psi_{xx}-\psi_{yy}\right)\right]\right\} \qquad (39)$$

The final equations describing our plasma are then

$$\frac{\partial n}{\partial t}-\frac{c}{B}\left(\nabla\psi\times\hat{e}_z\right)\cdot\nabla n = 0 \qquad (40)$$

$$\nabla\cdot\left[\left(\frac{c^2}{B^2}nm_i\right)\left(\frac{\partial}{\partial t}+\underset{\sim}{v}_{i1}\cdot\nabla+\nu_{in}\right)\nabla\psi\right] = \nabla\cdot\left[ne\frac{\nu_{in}}{\Omega_i}\underset{\sim}{U}_n\times\hat{e}_z\right]+\nabla\cdot\underset{\sim}{J}_Q+\nabla\cdot\underset{\sim}{J}_R \qquad (41)$$

Note that we have dropped the small term $\nabla\cdot[en\underset{\sim}{U}_n\nu_{in}^2/\Omega_i^2]$ in (41) to be consistent with Sperling and Glassman (1985). Note also that in the above equations, electron and ion temperature can change the plasma evolution only through η_1 and η_3. This is a result of our assumptions of no electron collisions, and of spatially uniform ion and electron temperatures.

III. STABILITY ANALYSIS

We consider the following equilibrium condition, which we denote by zero subscripts. We take n_0, and hence η_{10} and η_{30} to be functions of y only, and $\underset{\sim}{U}_n$ to be constant and in the y direction

$$\frac{\partial n_0}{\partial x} = \frac{\partial \eta_{10}}{\partial x} = \frac{\partial n_{30}}{\partial x} = 0 \tag{42}$$

$$\underset{\sim}{U}_n = U_n \hat{e}_y \tag{43}$$

This equilibrium configuration is characterized by

$$\psi_0 = 0 \tag{44}$$

$$\underset{\sim}{v}_{i10} = \frac{\nu_{in}}{\Omega_i} U_n \hat{e}_x \tag{45}$$

(The reader may have noticed that this appears not to be an equilibrium since $v_{i20} \neq 0$; this is a consequence of errors introduced in the Sperling and Glassman (1985) ordering. If the treatment had been the more accurate classical one, a true equilibrium would be obtained.)

We introduce the perturbation quantities

$$n(x,y,t) = n_0(y) + \tilde{n}(y,t)e^{ikx} \tag{46}$$

$$\psi(x,y,t) = \tilde{\psi}(y,t)e^{ikx}. \tag{47}$$

Linearizing (40) and (41) we find that

$$\frac{\partial \tilde{n}}{\partial t} + \frac{c}{B} ik\tilde{\psi} \frac{\partial n_0}{\partial y} = 0 \tag{48}$$

$$\frac{\partial}{\partial y}\left[\sigma_{p0}\frac{\partial\tilde{\psi}}{\partial y}\right] - \sigma_{p0}k^2\tilde{\psi} + \frac{\partial}{\partial y}\left[\frac{\sigma_{p0}}{\nu_{in}}\frac{\partial}{\partial y}\left(\frac{\partial}{\partial t}\tilde{\psi} + ik\frac{\nu_{in}}{\Omega_i}U_n\tilde{\psi}\right)\right]$$

$$- \frac{\sigma_{p0}}{\nu_{in}}k^2\left(\frac{\partial}{\partial t}\tilde{\psi} + ik\frac{\nu_{in}}{\Omega_i}U_n\tilde{\psi}\right) = \frac{\nu_{in}}{\Omega_i}U_n e^{ik\tilde{n}} + \nabla\cdot\tilde{J}_Q + \nabla\cdot\tilde{J}_R \tag{49}$$

where

$$\sigma_{p0} = \frac{n_0 ce}{B}\frac{\nu_{in}}{\Omega_i} \tag{50}$$

$$\nabla\cdot\tilde{J}_Q = \frac{c^2}{B^2}\left[2\eta_{10y}\left(\tilde{\psi}_{xxy} + \tilde{\psi}_{yyy}\right) + \eta_{10yy}\left(\tilde{\psi}_{yy} - \tilde{\psi}_{xx}\right) + \eta_{10}\left(\tilde{\psi}_{xxxx} + 2\tilde{\psi}_{xxyy} + \tilde{\psi}_{yyyy}\right)\right] \tag{51}$$

$$\nabla\cdot\tilde{J}_R = \frac{c^2}{B^2}\left[-2\eta_{30y}\left(\tilde{\psi}_{xxx} + \tilde{\psi}_{xyy}\right) - 2\eta_{30yy}\tilde{\psi}_{xy}\right] \tag{52}$$

The usual approach to stability analysis is to use the following substitutions

$$\frac{\partial}{\partial x} \to ik \tag{53}$$

$$\frac{\partial}{\partial t} \to i\omega \tag{54}$$

in (48) and (49); thereby eliminating either $\tilde{n}$ or $\tilde{\psi}$, and obtaining an eigenvalue problem for a fourth order differential operator. This is the approach used in Sperling and Glassman (1985), which yielded an equation they solved only in the long wavelength limit. Here we solve the problem with no approximations by treating equations (48) - (52) as an initial

value problem, using a random seed perturbation. All derivatives, including temporal derivatives, are discretized using standard finite differences. Assuming that there is a fastest-growing eigenmode, this mode will emerge from the noise, eventually reaching a point where, on a relative scale, it is virtually the only mode left in the system. At this point both the fastest growing eigenmode and its complex eigenfrequency have been isolated. Before we move on to these results, let us look at two limiting cases for which solutions are somewhat easier to obtain: the asymptotic cases of long- and short-wavelength limits.

IV. ASYMPTOTIC RESULTS: SHORT- AND LONG-WAVELENGTH LIMITS

A. Short-Wavelength Limit

We consider first the short-wavelength limit $kL \gg 1$, where $L \equiv n_0(\partial n_0/\partial y)^{-1}$ is the gradient scale length. We take the perturbation quantities in (48) - (52) to be invariant in y. Making the substitutions (53) and (54), (48) and (49) become

$$i\omega\tilde{n} = -\frac{c}{B}\, ik\, \frac{\partial n_0}{\partial y}\, \tilde{\psi} \tag{55}$$

$$-\sigma_{p0}k^2\,\tilde{\psi} - I\,\frac{\sigma_{p0}}{\nu_{in}}\,k^2\left(i\omega + ik\,\frac{\nu_{in}}{\Omega_i}\,U_n\right)\tilde{\psi} = \frac{\nu_{in}}{\Omega_i}\,U_n eik\tilde{n}$$

$$+\frac{c^2}{B^2}\left[\eta_{10yy}k^2 + \eta_{10}k^4 + 2\eta_{30y}ik^3\right]\tilde{\psi} \tag{56}$$

where I is a flag on ion inertia (I = 1 to retain ion inertia, I = 0 to neglect it). Solving (56) for $\tilde{n}$, and substituting into (55) we obtain

$$i\omega\left[\frac{\Omega_i}{i\nu_{in}U_n ek}\right]\left[-\sigma_{p0}k^2 - I\frac{\sigma_{p0}}{\nu_{in}}k^2\left(i\omega + ik\frac{\nu_{in}}{\Omega_i}U_n\right)\right.$$

$$\left. - \frac{c^2}{B^2}\left(\eta_{10yy}k^2 + \eta_{10}k^4 + 2\eta_{30y}ik^3\right)\right] = -\frac{c}{B}ik\frac{\partial n_0}{\partial y} \qquad (57)$$

Noting that $(\Omega_i/\nu_{in}e)\,\sigma_{p0} = (c/B)n_0$, dividing through by this quantity, and multiplying by iU_nk, we obtain

$$i\omega\left[-k^2 - Ik^2\nu_{in}^{-1}\left(i\omega + ik\frac{\nu_{in}}{\Omega_i}U_n\right) - \frac{c^2}{B^2}\sigma_{p0}^{-1}\left(\eta_{10yy}k^2\right.\right.$$

$$\left.\left. + \eta_{10}k^4 + 2\eta_{30y}ik^3\right)\right] = k^2\frac{1}{n_0}\frac{\partial n_0}{\partial y}U_n \qquad (58)$$

This is simply a quadratic equation for ω of the form

$$A\omega^2 + B\omega + C = 0 \qquad (59)$$

with

$$A = I\nu_{in}^{-1} \qquad (60)$$

$$B = -i + I\nu_{in}^{-1}k\frac{\nu_{in}}{\Omega_i}U_n - i\frac{c^2}{B^2}\sigma_{p0}^{-1}\left(\eta_{10yy} + \eta_{10}k^2 + 2\eta_{30y}ik\right) \qquad (61)$$

$$C = -\frac{1}{n_0}\frac{\partial n_0}{\partial y}U_n \qquad (62)$$

An instructive look at how the η_1 and η_3 terms might contribute to the stabilization of the gradient drift instability is obtained by first dropping ion inertia (setting I = 0), and looking at the collisional limit. Then

$$\omega = iU_n \frac{1}{n_0} \frac{\partial n_0}{\partial y} \left[1 + \frac{c^2}{B^2} \sigma_{p0}^{-1} \left(\eta_{10yy} + \eta_{10} k^2 + 2\eta_{30y} ik\right)\right]^{-1} \tag{63}$$

The growth rate γ is given by minus the imaginary part of ω:

$$\gamma = \frac{- U_n \frac{1}{n_0} \frac{\partial n_0}{\partial y} \left[1 + \frac{c^2}{B^2} \sigma_{p0}^{-1} \left(\eta_{10yy} + \eta_{10} k^2\right)\right]}{\left[1 + \frac{c^2}{B^2} \sigma_{p0}^{-1} \left(\eta_{10yy} + \eta_{10} k^2\right)\right]^2 + \left[2 \frac{c^2}{B^2} \sigma_{p0}^{-1} \eta_{30y} k\right]^2} \tag{64}$$

One can note that the η_3 terms always act to stabilize the collisional short-wavelength limit, but that the η_1 terms need not always be similarly stabilizing.

In the limit of small ν_{ii}/Ω_i (low electron density) (25) - (31) yield

$$\eta_1 = 1.2 \frac{n_i k_B T_i}{\nu_{ii}} \frac{\nu_{ii}^2}{4\Omega_i^2} = 0.3 \, n k_B T_i \frac{\nu_{ii}}{\Omega_i^2} \tag{65}$$

$$\eta_3 = \frac{n_i k_B T_i}{\nu_{ii}} \frac{\nu_{ii}}{2\Omega_i} = \frac{1}{2} n k_B T_i / \Omega_i \tag{66}$$

Since the mean thermal ion gyroradius ρ_i is given by

$$\rho_i^2 = \frac{k_B T_i}{m_i \Omega_i^2} \tag{67}$$

and since

$$\frac{c^2}{B^2} \sigma_p^{-1} = \frac{\Omega_i}{\nu_{in}} \frac{B}{nce} \frac{c^2}{B^2} = \frac{\Omega_i}{\nu_{in}} \frac{c}{Bne} \tag{68}$$

then

$$2\,\frac{c^2}{B^2}\,\sigma_{p0}^{-1}\,\eta_{30y} = 2\,\frac{\Omega_i}{\nu_{in}}\,\frac{c}{Bn_0e}\,\frac{\partial}{\partial y}\left(\frac{1}{2}\,n_0\rho_i^{\,2}\,m_i\Omega_i\right) = \frac{\Omega_i}{\nu_{in}}\,\rho_i^{\,2}\,\frac{1}{n_0}\,\frac{\partial n_0}{\partial y} \tag{69}$$

Thus, if we let $\nu_{ii}/\Omega_i \to 0$ we get

$$\gamma = -\,\frac{U_0/L}{1 + \Omega_i^{\,2}k^2\rho_i^{\,4}/L^2\nu_{in}^{\,2}} \tag{70}$$

where we have defined

$$L^{-1} = \frac{1}{n_0}\left(\frac{\partial n_0}{\partial y}\right) \tag{71}$$

Recall that this expression is valid for the collisional, short wavelength limit in the low electron density regime ($\nu_{ii}/\Omega_i \to 0$). Huba and Zalesak (1985) have shown that the long wavelength limit ($kL \ll 1$) can often be obtained from the short wavelength limit by making the simple substitution $1/L \to k(M - 1)/(M + 1)$, yielding

$$\gamma = -\,\frac{M - 1}{M + 1}\,kU_0\left[1 + \left(\frac{M - 1}{M + 1}\right)^2 k^4\hat{\rho}_i^{\,4}\right]^{-1} \tag{72}$$

$$\hat{\rho}_i \equiv \sqrt{\frac{\Omega_i}{\nu_{in}}}\,\rho_i \tag{73}$$

Here M is the ratio of the electron densities on either side of the discontinuity. This collisional long wavelength limit has been obtained rigorously by Sperling and Glassman (1985). It is obvious that η_3 effects

will strongly damp the instability whenever

$$k\rho_i \geq \left(\nu_{in}/\Omega_i\right)^{1/2} \tag{74}$$

in this short wavelength collisional limit.

B. Long-Wavelength Limit

By the long-wavelength limit ($kL \ll 1$) we mean the solution to the linearized equations (47) - (52) with

$$n_0(y) = \begin{cases} n_< & ; \quad y < y_0 \\ n_> = Mn_< & ; \quad y \geq y_0 \end{cases} \tag{75}$$

The concept of the long wavelength limit within the context of the stability of ionospheric ion clouds was introduced by Huba and Zalesak (1983), and has become a standard analysis tool in the field. The logic behind the use of the long wavelength limit rather than the older short wavelength limit is that the clouds in question tend to steepen considerably before structuring, making the analysis of a very steep density jump a more plausible approximation than that of a smooth exponential profile. One of the primary purposes of this paper is to compare the long wavelength results with those of continuous, more realistic profiles, in order to ascertain the applicability of the long wavelength approximation for realistic clouds.

The equations describing the long-wavelength limit are derived in Sperling and Glassman (1985), and we shall not repeat that analysis here. Suffice it to say that the equations are easier to solve than the full set

(48) - (52), and that we have developed software which reproduces the Sperling and Glassman (1985) results, for comparison purposes.

V. NUMERICAL RESULTS

A. Long-Wavelength Comparisons

All of the results we report here are for an atomic oxygen plasma ($m_i = 16.0\ m_p$), $T_i = 0.1$ eV, and use the following profile for n_0

$$n_0 = n_< \left[1 + \frac{M - 1}{2}\left(1 + \tanh\left(\frac{y - y_0}{L_0}\right)\right)\right] \tag{76}$$

$n_< = 10^5\ cm^{-3}$, $M = 1000$, and y_0 is the center of our computational domain in all of our problems. Three different values of ν_{in} are used, meant to span a wide range of altitudes: $\nu_{in} = 10\ sec^{-1}$ (160 km altitude), $\nu_{in} = 1\ sec^{-1}$ (250 km altitude) and $\nu_{in} = 0.1\ sec^{-1}$ (500 km altitude). The neutral wind velocity U_n was taken to be $U_n = -100$ m/s in all cases.

Since our primary goal in this paper is to examine the applicability of long wavelength asymptotic theory to realistic situations with finite gradients, we shall compare calculations in which only L_0 changes in (76). Note that the long wavelength limit corresponds to $L_0 \to 0$.

Figure 1 shows results for the case $\nu_{in} = 0.1\ sec^{-1}$. Displayed is a plot of γ vs. k for $L_0 = 0$, 20 m, 100 m, 250 m, 500 m and 1000 m. Figures 2 and 3 show results for the cases $\nu_{in} = 1\ sec^{-1}$ and $10\ sec^{-1}$, respectively. The solid curve is for $L_0 = 0$ m and is based on the exact asymptotic dispersion equation; the curves labelled A, B, C, D, and E are for $L_0 = 20$ m, 100 m, 250 m, 500 m, and 1000 m, respectively, and are based on numerical solutions to (48) - (52). Missing data corresponds to cases where convergence to an unambiguous fastest-growing eigenmode could

not be obtained with the present code.

It is obvious from Fig. 1-3 that the value of L does substantially affect the growth rate γ. There are two equally important aspects of this issue which bear on our discussion here. First, there is the absolute effect of L on γ, all other parameters being held fixed. From a look at Figs. 1-3, the reader can see that the growth rate for a finite L is always smaller than that for $L = 0$, and that, in general, increasing L decreases γ, all other parameters being held fixed. The degree to which this reduction takes place does seem to be a function of the ion-neutral collision frequency, with the effect being largest for large ν_{in}, and smallest for small ν_{in}.

The second, perhaps more important, aspect of the effects of finite L, is the degree to which L affects the k for which $\gamma(k)$ maximizes, k_{max}. We shall also be interested in the degree to which the curve $\gamma(k)$ is sharply peaked about k_{max}. Both of these issues are highly relevant to the freezing models which have been proposed [Glassman and Sperling, 1987, Zalesak et al., 1988]. Again referring to Figs. 1-3, the reader will note that the effect of increasing L is to make k_{max} smaller. Since k_{max} is in fact the freezing scale for many freezing models, this is a very important effect. The effect of increasing L on the degree to which the curve $\gamma(k)$ is sharply peaked is more difficult to describe. As L is increased from 0 to 20 m. the curve becomes more sharply peaked, which is a desirable attribute for the freezing model of Zalesak et al. (1988). However, it is clear that as L is further increased, the curve eventually becomes less sharply peaked with increasing L. Thus larger values of L may make some freezing models which depend on $\gamma(k)$ being sharply peaked about k_{max} less desirable.

B. Short-Wavelength Comparisons

Having compared true growth rates with long-wavelength asymptotic results, and found the asymptotic results lacking to some degree, especially at short wavelengths, it is natural to assume that perhaps the short-wavelength limit, (58) - (62), might be a fairly good approximation in this regime. However, it is not clear how to apply (58) - (62) to our profile (76), since the spatial derivatives of n_0, n_1, and n_3 are functions of y. We have chosen to evaluate $\gamma(y)$ as given by (58) - (62) at every point in y of the profile (76), and to take

$$\gamma_{SWL} = \max \gamma(y), \; -\infty \leq y \leq +\infty \tag{77}$$

In Fig. 4 we compare γ_{SWL} with the numerically computed solution for the case L = 1000 m, for all three values of ν_{in} (0.1, 1.0, 10.0 s^{-1}). Note that there is excellent agreement for kL > 10, but that there can be substantial disagreement for kL ~ 1. In Figs. 5-7, we show the same comparison for the cases L = 500 m, L = 250 m, and L = 100 m, respectively. Again, note that kL must be considerably larger than 1 to obtain good agreement between the curves, and that there are substantial errors in γ_{SLW} for kL ~ 1.

We introduce some bit of mystery here. The reader may note by comparing curve D (ν_{in} = 0.1) of Fig. 4-7 that γ_{SWL} seems to be asymptoting to a γ <u>independent of L</u> for large k in this highly collisionless case. We do not as yet have an explanation for this phenomenon.

C. Relative Effectiveness of η_1 and η_3

Finally in this section we wish to give the reader a feeling for the relative effectiveness of the η_1 and η_3 terms on the growth rate of γ. In Fig. 8 we show a plot of γ vs. k for the case L_0 = 20 m, and all three values of ν_{in} (0.1, 1.0, and 10.0 s^{-1}), but with $\eta_1 = \eta_3 = 0$. In Fig. 9 we show an identical plot, but for $\eta_3 = 0$ (with η_1 taking on its actual values). In Fig. 10 we show a plot with both η_1 and η_3 set to actual values. Note that the achievement of peaked values of $\gamma(k)$ came only when η_3 terms were taken into account, making it the primary piece of relevant physics for freezing models based on a k_{max}. This is verified by Fig. 11, where we show a similar plot with $\eta_1 = 0$, and only η_3 terms "turned on". We do not as yet have an explanation for the turning up of $\gamma(k)$ for ν_{in} = 0.1 at very large k in Fig. 11. (We find that making L smaller makes the turning up vanish, in agreement with the true L = 0 result.)

VI. CONCLUSIONS

We have examined the role of magnetic and collisional viscosity (the η_3 and η_1 terms, respectively, in Braginskii's stress tensor) on the $\underline{E} \times \underline{B}$ gradient drift instability, paying special attention to the degree to which the gradient scale length L affects curves of $\gamma(k)$. In particular, we wished to know the size of the error introduced by assuming L to be zero (the long-wavelength limit). We conclude that in addition to affecting γ in an absolute sense, increasing L to a finite value decreases k_{max}, the k for which $\gamma(k)$ maximizes, and also affects width of the peak of $\gamma(k)$ at k_{max}. Both of these results have large effects on freezing models. We have also examined the applicability of the old short-wavelength limit, or local analysis, and found that significant errors would be made unless kL is very large (~ 10).

ACKNOWLEDGMENT

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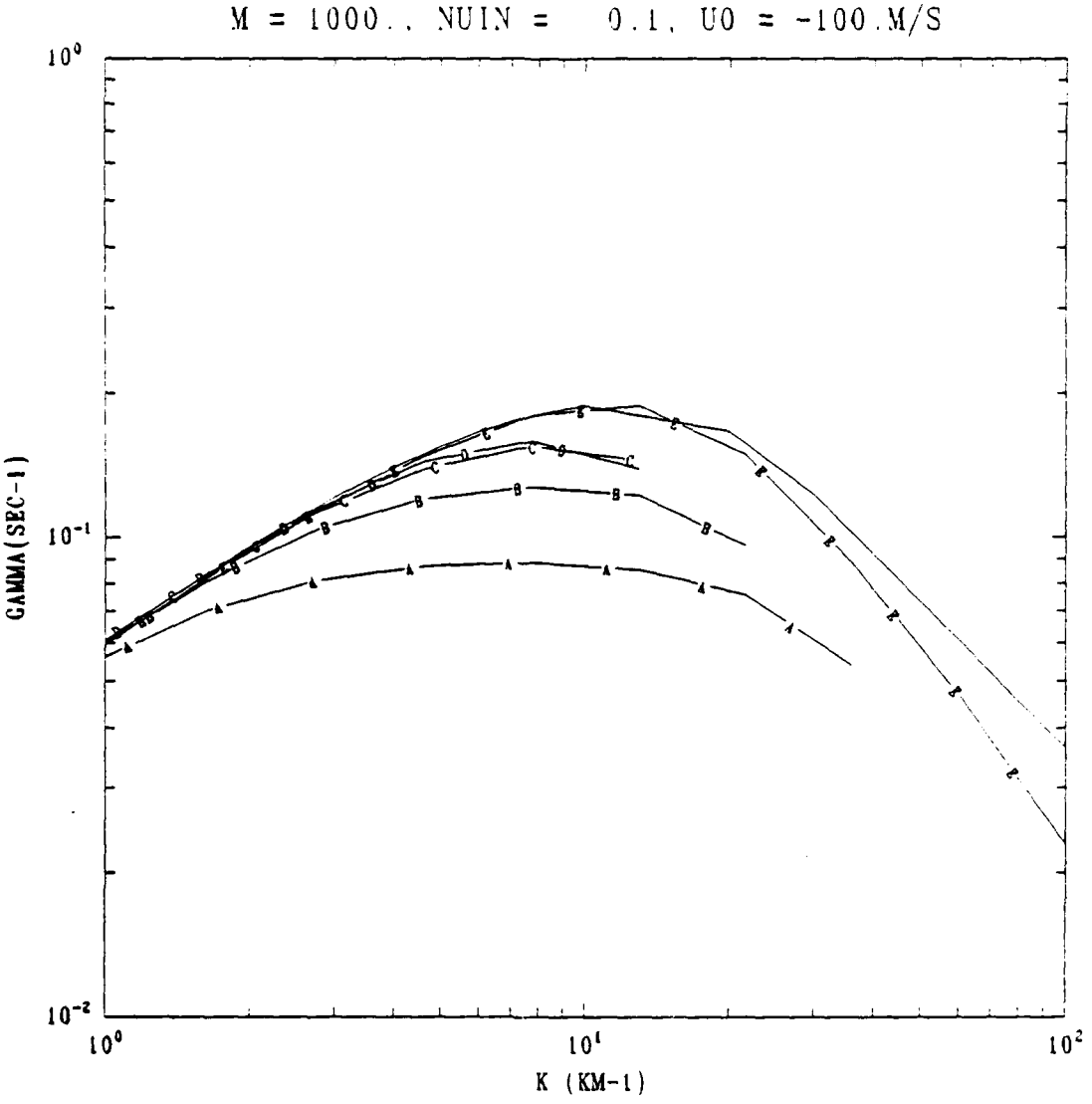

Fig. 1 Plot of γ vs. k for the case of $n_< = 10^5$, M = 1000, U_n = - 100 m/s, using a hyperbolic tangent density profile and $\nu_{in} = 0.1\ s^{-1}$. Curves A, B, C, D, and E refer to L = 1000 m, 500 m, 250 m, 100 m, and 20 m, respectively. Missing data corresponds to cases for which our code could not converge to an unambiguous fastest-growing eigenmode. The solid unlabeled curve is the long-wavelength limit (L = 0).

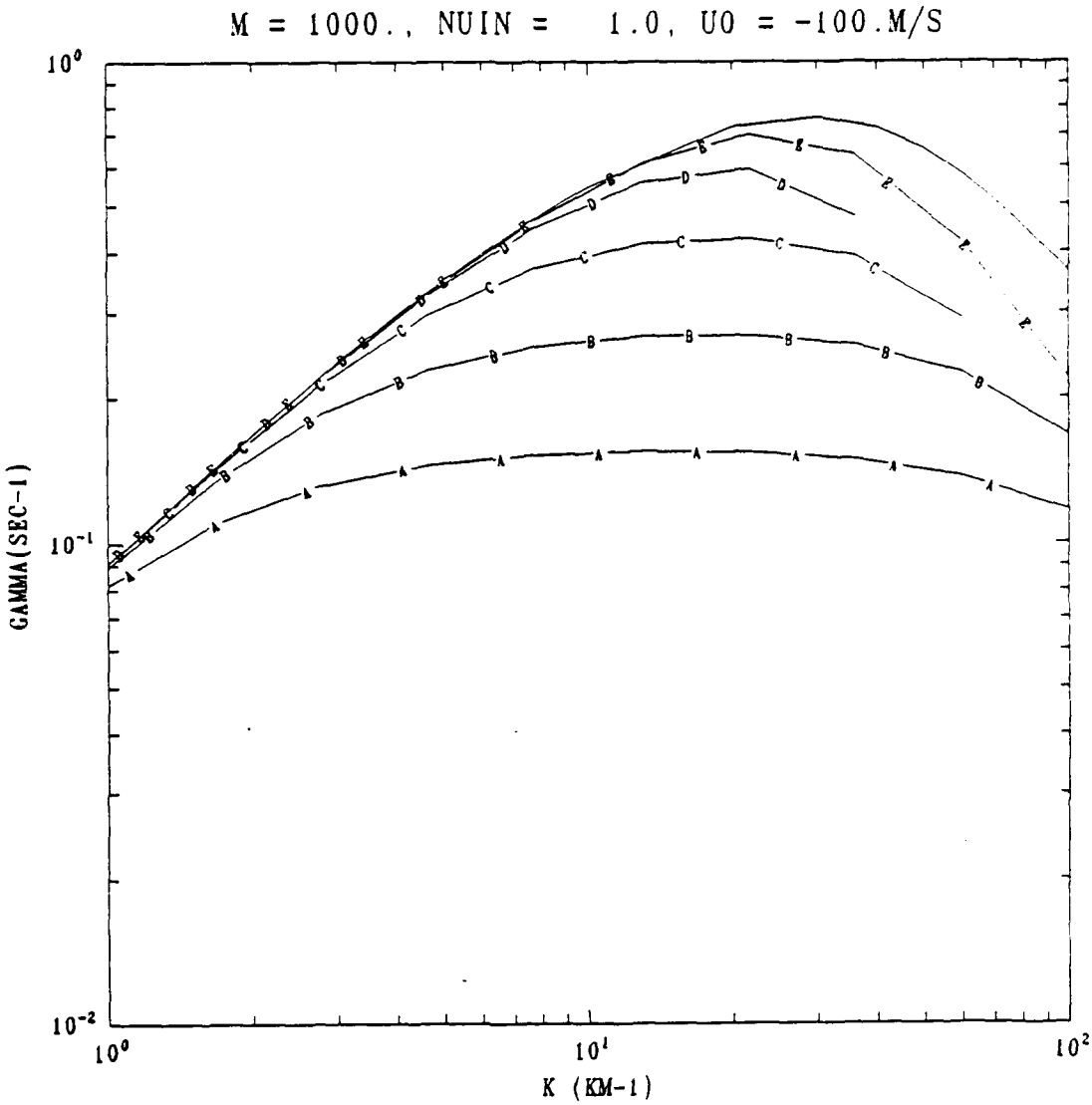

Fig. 2 As in Fig. 1, but for $\nu_{in} = 1.0\ s^{-1}$.

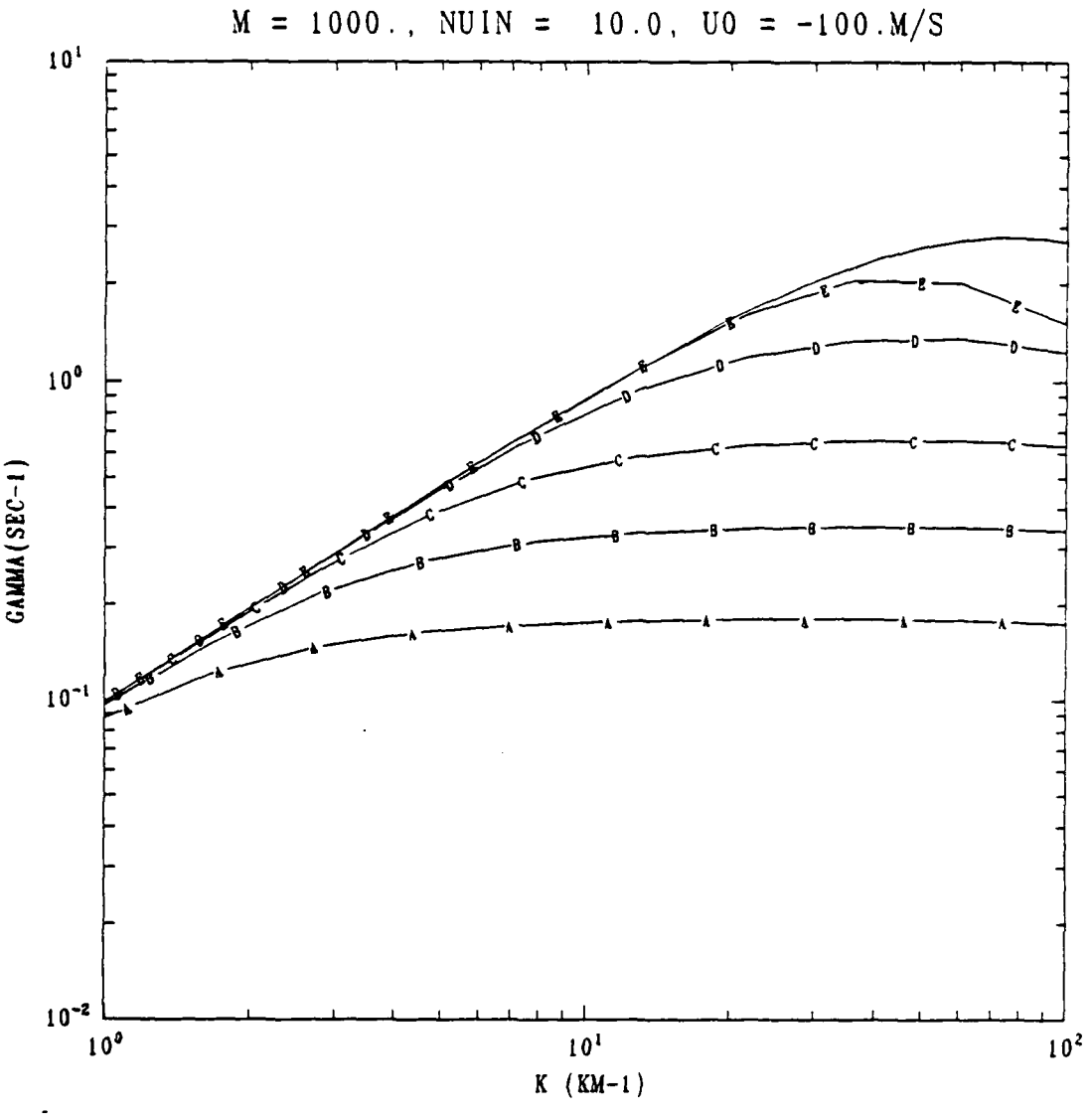

Fig. 3 As in Fig. 1, but for ν_{in} = 10.0 s^{-1}.

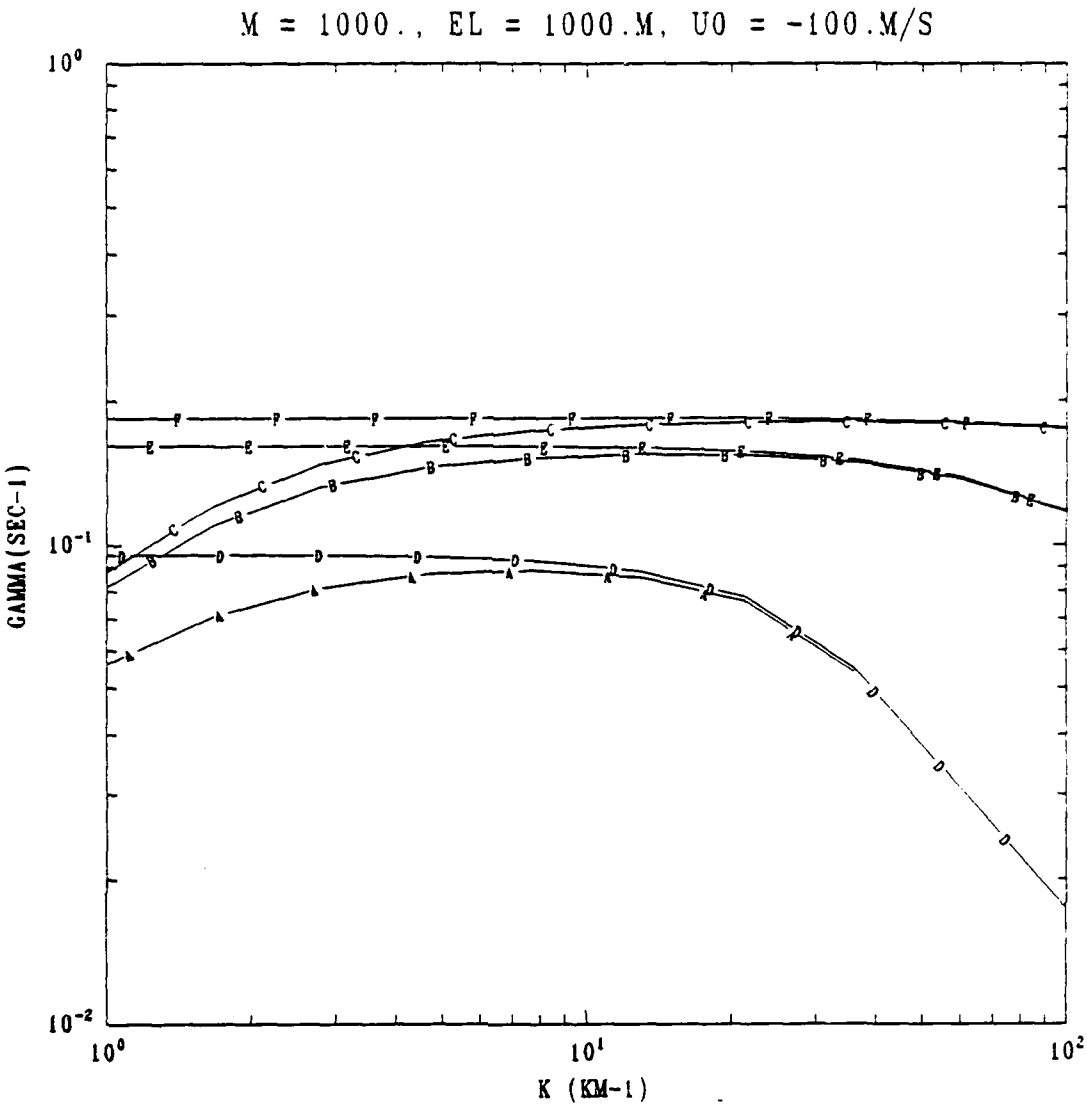

Fig. 4 Plots of γ vs. k for the case of $n_< = 10^5$, M = 1000, U_n = - 100 m/s, using a hyperbolic tangent density profile, and L = 1000 m. Curves A, B, and C refer to ν_{in} = 0.1, 1.0, and 10.0 s^{-1}, respectively. Curves D, E, and F refer again to ν_{in} = 0.1, 1.0, and 10.0 s^{-1}, respectively, but are generated using the short-wavelength asymptotic limit (77), (58)-(62).

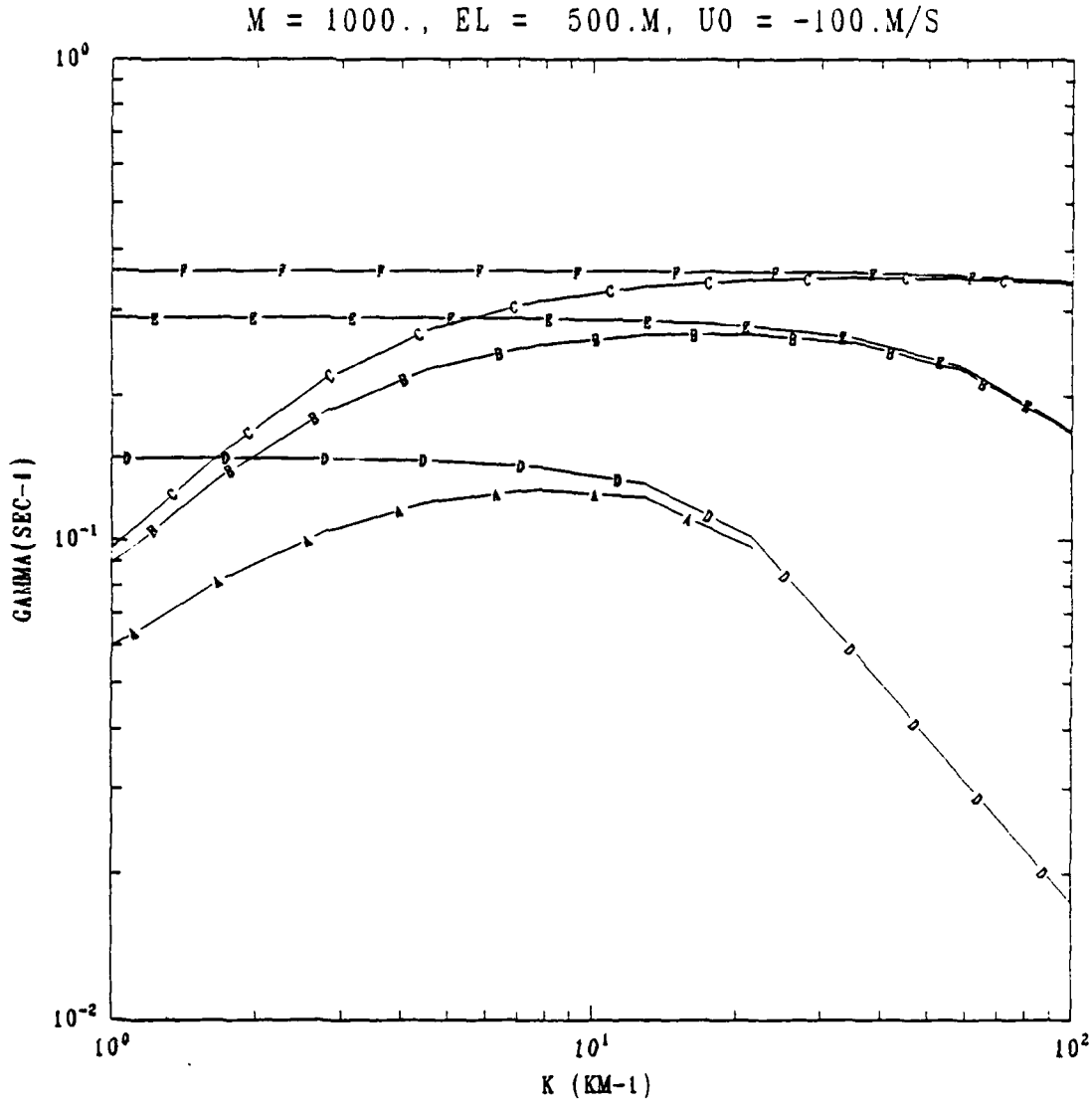

Fig. 5 As in Fig. 4, but for L = 500 m.

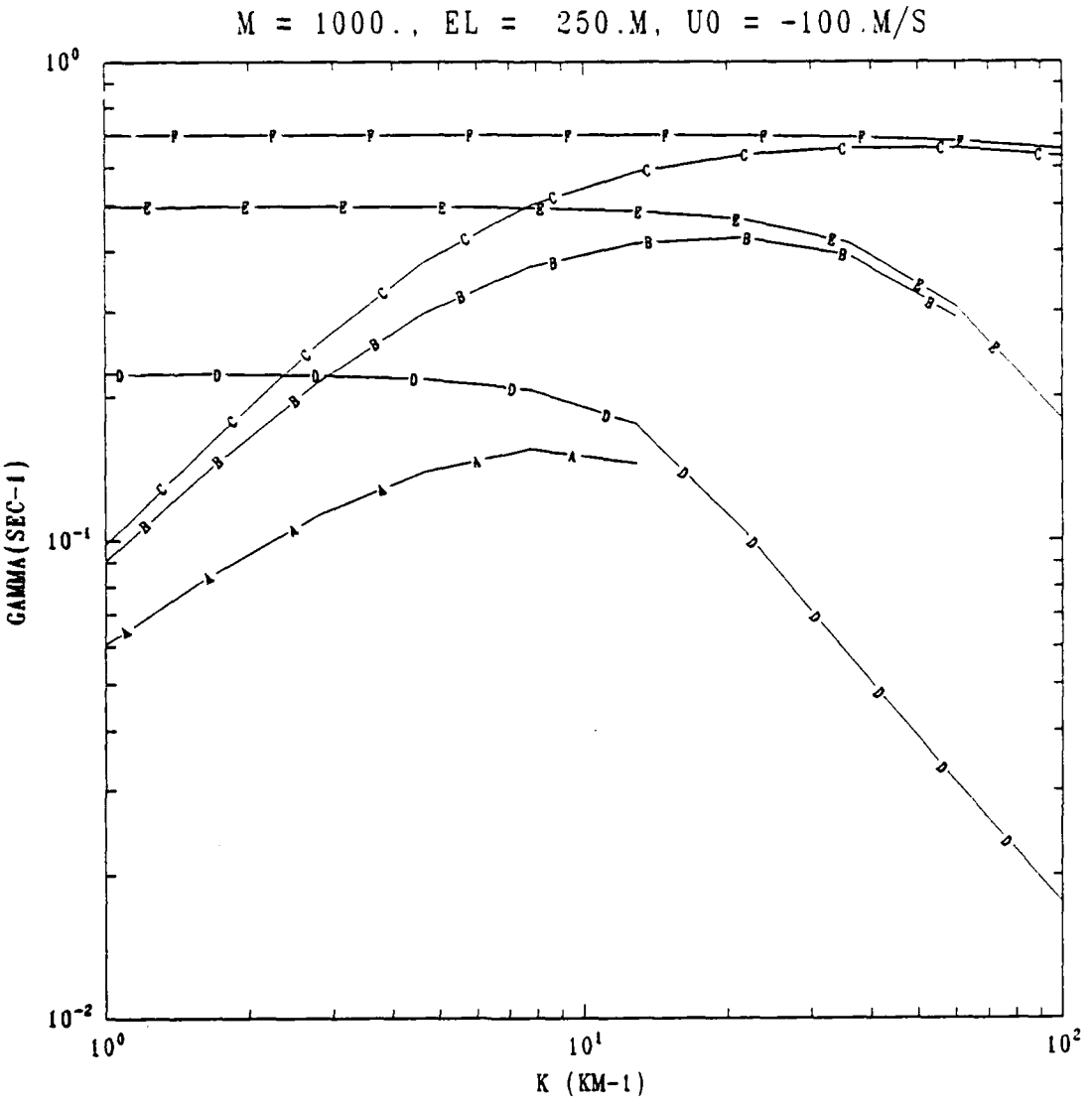

Fig. 6 As in Fig. 4, but for L = 250 m.

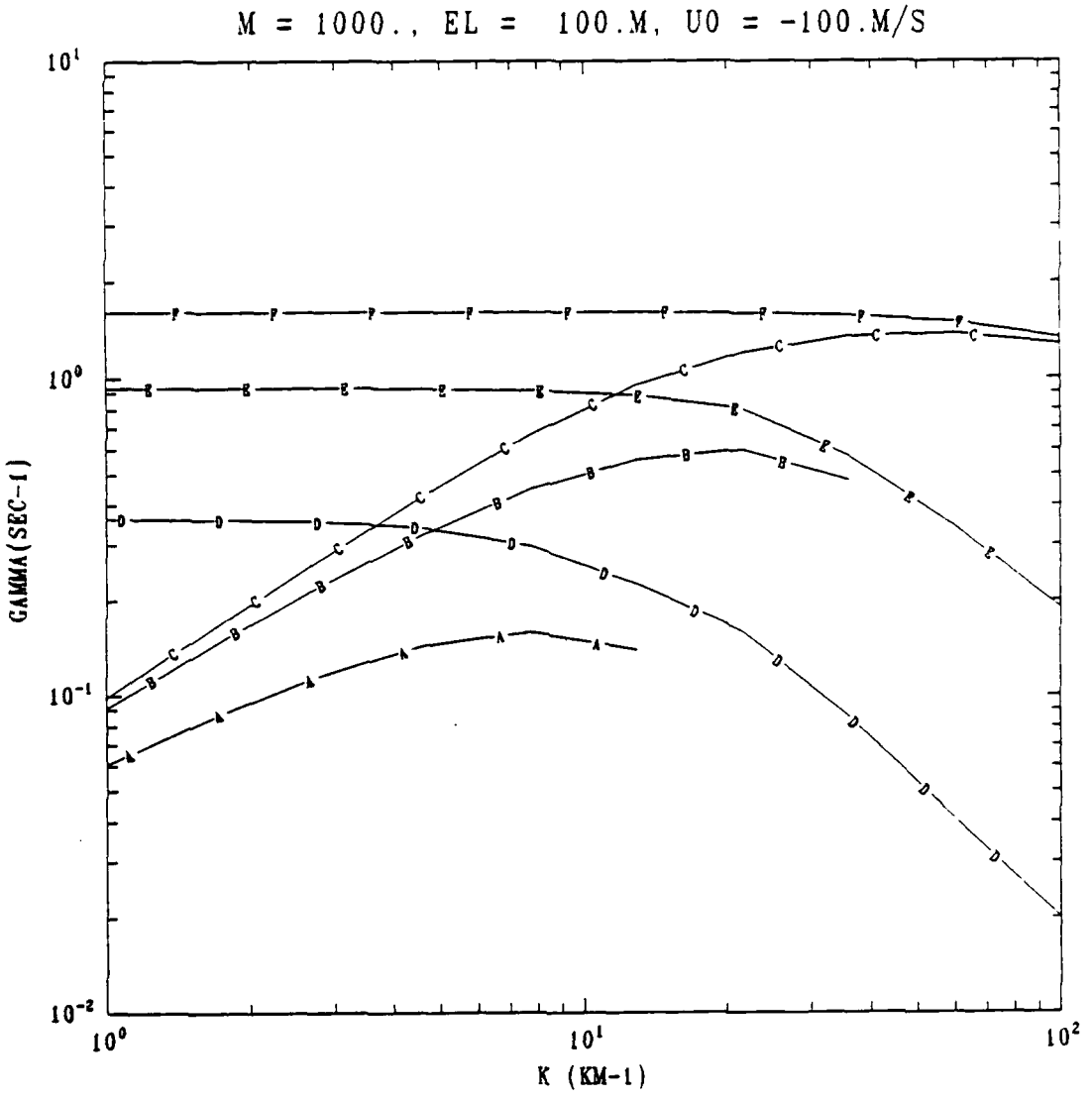

Fig. 7 As in Fig. 4, but for L = 100 m.

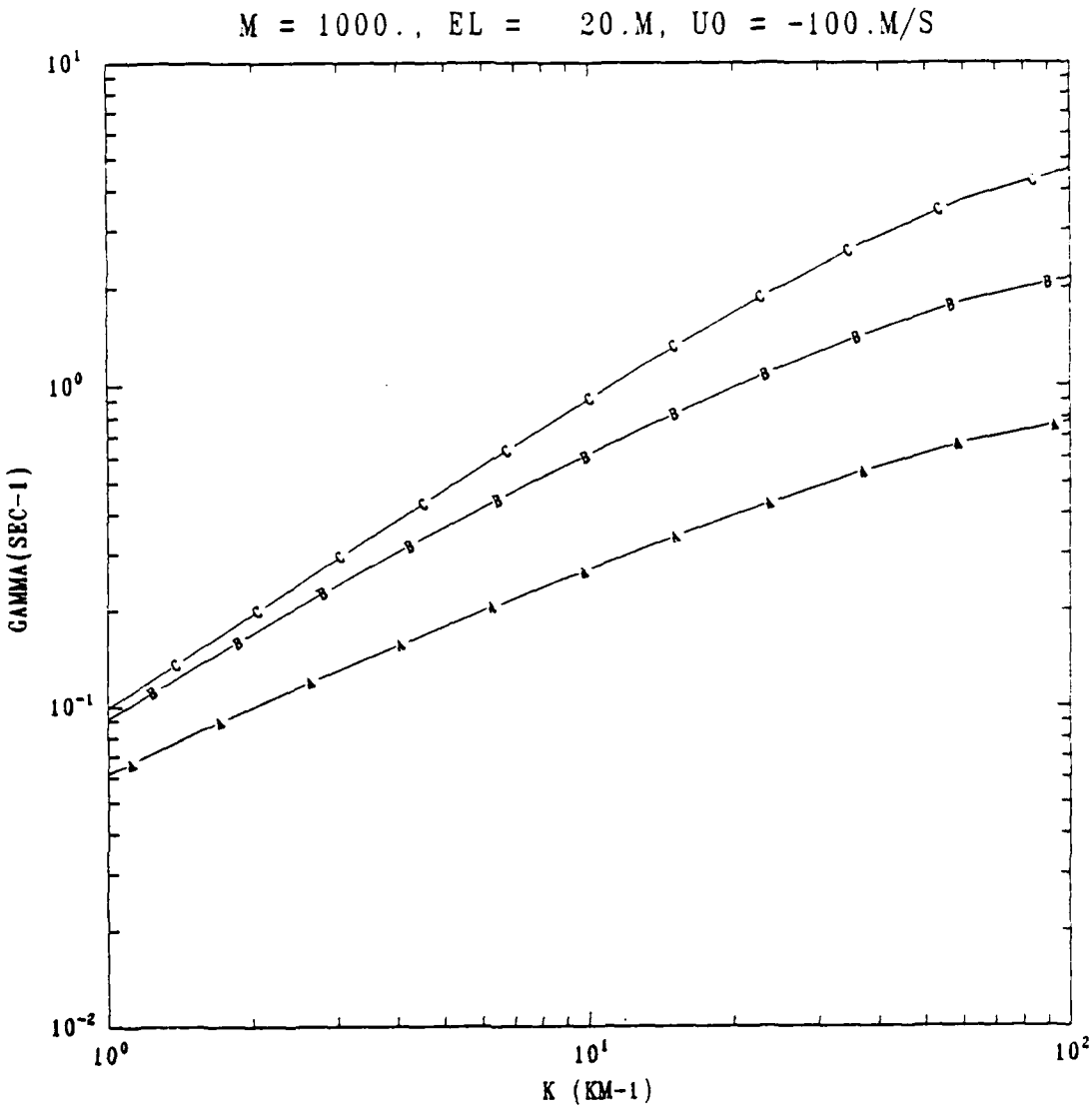

Fig. 8 Plots of γ vs. k for the case of $n_< = 10^5$, M = 1000, U_n = - 100 m/s, using a hyperbolic tangent density profile, and L = 20 m. Both the collisional viscosity η_1 and the magnetic viscosity η_3 have been neglected. Curves A, B, and C refer to ν_{in} = 0.1, 1.0, and 10.0 s^{-1}, respectively.

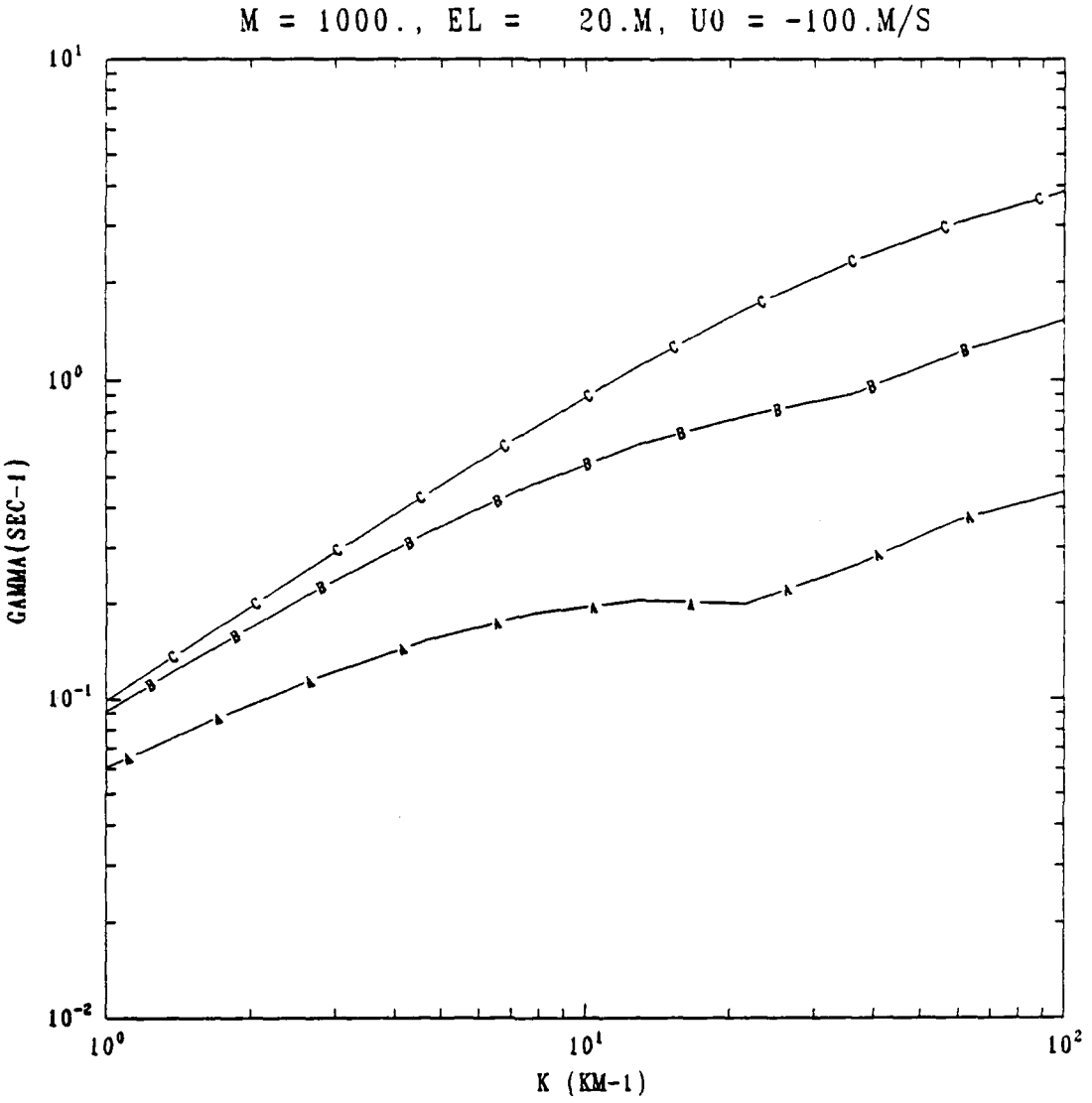

Fig. 9 As in Fig. 8 but only the magnetic viscosity η_3 has been neglected.

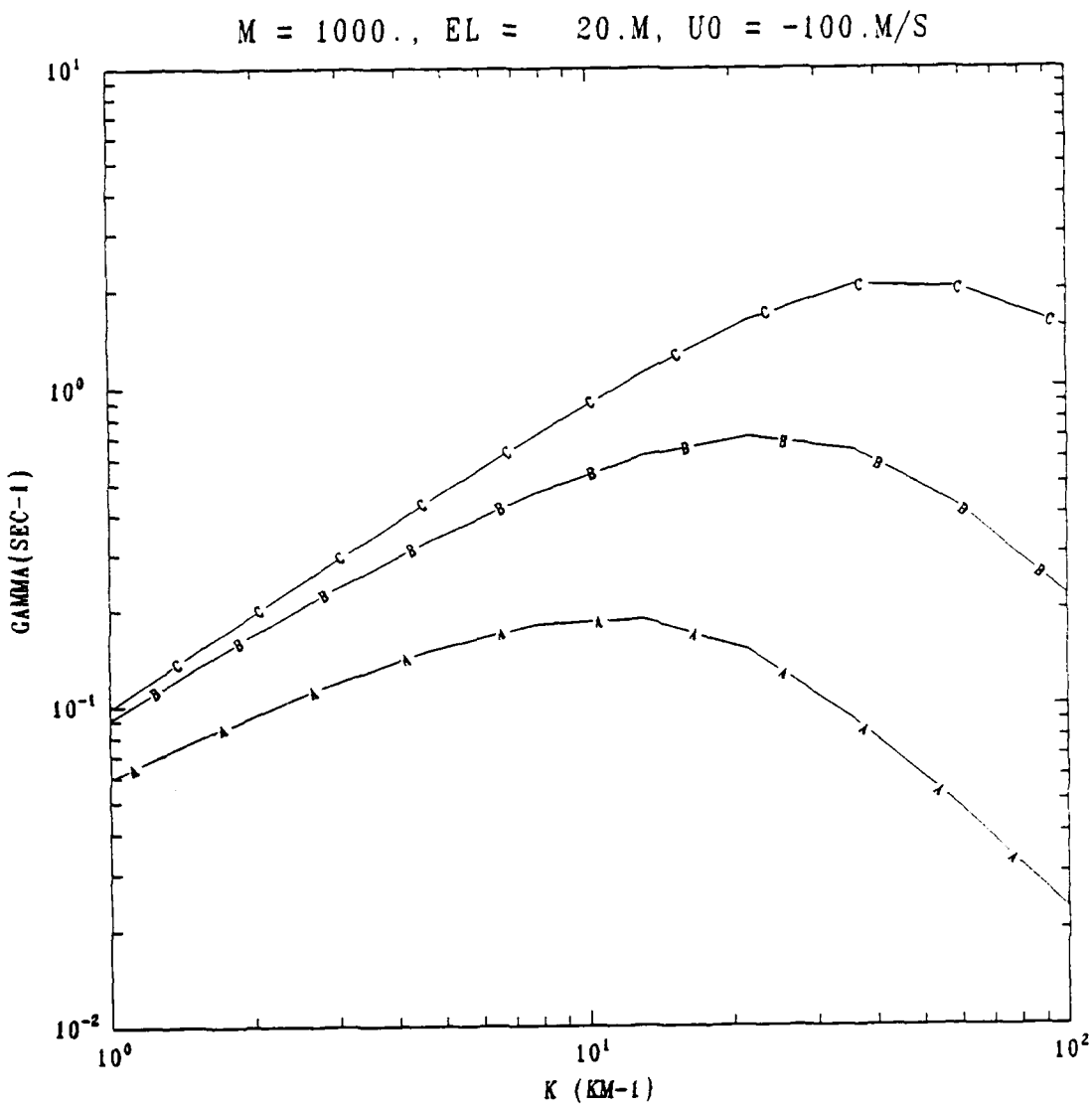

Fig. 10 As in Fig. 8 but neither the collisional viscosity nor the magnetic viscosity has been neglected.

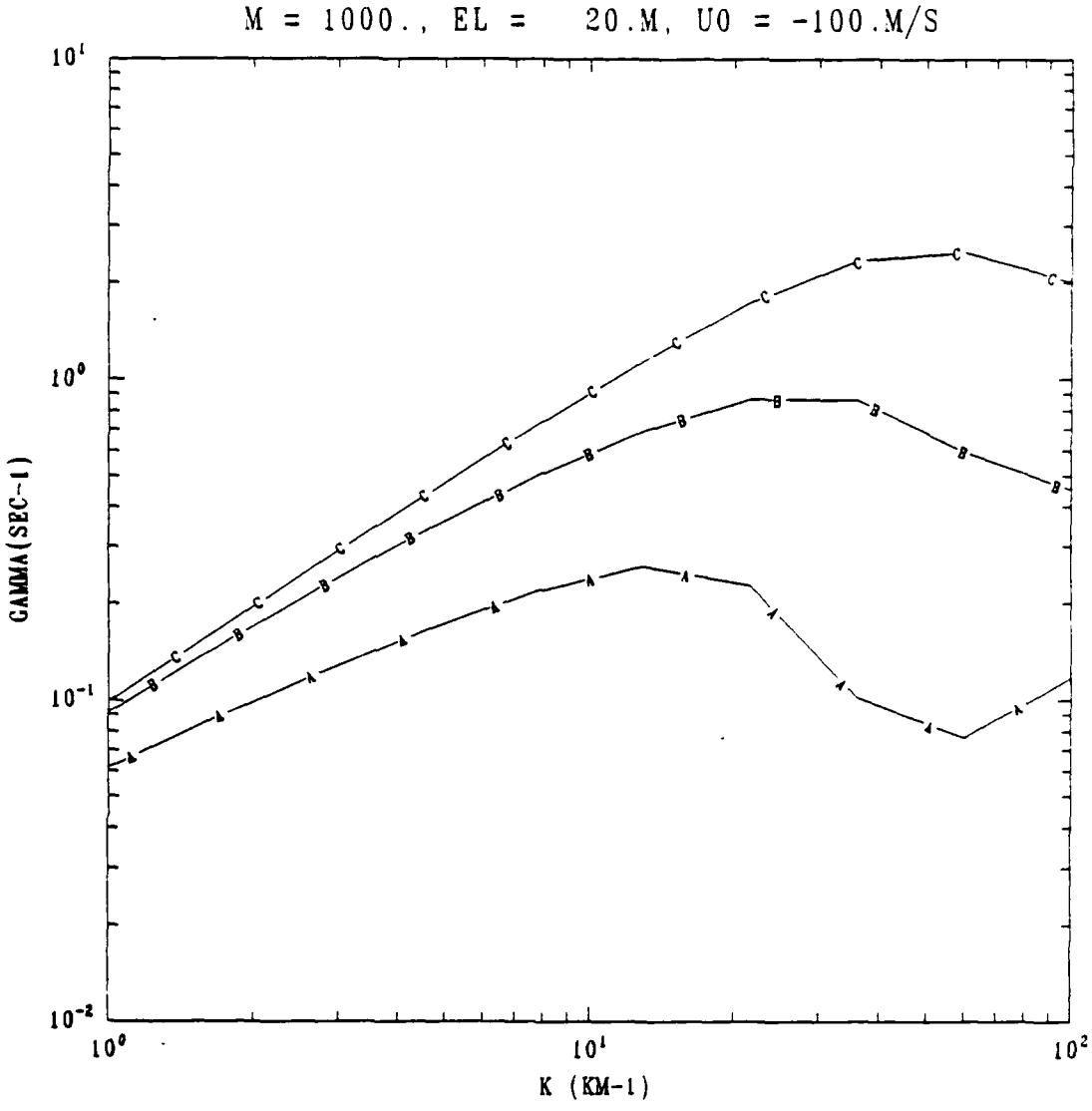

Fig. 11 As in Fig. 8, but only the collisional viscosity η_1 has been neglected.

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